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# VALIDATION OF A PSEUDO-SOUND THEORY FOR THE PRESSURE-DILATATION IN DNS OF COMPRESSIBLE TURBULENCE

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**Abstract.** The results of an asymptotic theory for statistical closures for compressible turbulence are explored and validated with the direct numerical simulation of the isotropic decay and the homogeneous shear. An excellent collapse of the data is seen. The slow portion is found to scale, as predicted by the theory, with the quantity  $M_t^2$  and  $\varepsilon_s$ . The rapid portion has an unambiguous scaling with  $\alpha^2 M_t^2 \varepsilon_s [\frac{P_k}{\varepsilon} - 1] (\frac{Sk}{\varepsilon_s})^2$ . Implicit in the scaling is a dependence, as has been noted by others, on the gradient Mach number. A new feature of the effects of compressibility, that of the Kolmogorov scaling coefficient,  $\alpha$ , is discussed. It is suggested that  $\alpha$  may contain flow specific physics associated with large scales that might provide further insight into the structural effects of compressibility.

**Key words.** pressure-dilatation, compressible turbulence, turbulence modeling

**Subject classification.** Fluid Mechanics

**1. Introduction.** The validation, using DNS data, of the scaling predicted by an analytical development for the pressure dilatation,  $\langle pd \rangle$ , appearing in the single point closures for compressible turbulence, given in Ristorcelli (1995, 1997), is the subject of this article. The analysis is relevant to shear flows with negligible bulk dilatation and low  $M_t^2$ . These restrictions are satisfied in a wide number of flows ranging from simple shear layers of theoretical interest, Papamoschou and Roshko (1988), to the complex shear layers associated with supersonic mixing enhancement, Gutmark *et al.* (1995). In most of these supersonic shear layers a Mach number based on the fluctuating velocity of the fluid particle is small. A Mach 4 mean flow with a turbulence intensity of 8 *per cent* has a turbulent Mach number of  $M_t = 0.32$ . The square of this turbulent Mach number, the appropriate perturbation expansion parameter arising from the Navier-Stokes equations,  $M_t^2 \sim 0.1$ , is small. The existence of this small parameter,  $M_t^2$ , allows some analytical results.

The article is primarily a study of the analysis of Ristorcelli (1995, 1997) in the light of recent DNS results of Blaisdell. The representations given in Ristorcelli (1995, 1997) were obtained using scaling arguments about the effects of compressibility, a singular perturbation idea and the methods of statistical fluid mechanics. While the results are expressed in the context of a statistical turbulence closure, they provide, with few phenomenological assumptions, an interesting and clear physical model for the scalar effects of compressibility. For a homogeneous turbulence with quasi-normal isotropic large scales the expressions derived are – in the small turbulent Mach number *squared* limit – exact. The analytical results, which are a rigorous consequence of the low  $M_t^2$  assumptions and do not contain any unspecified empirical coefficients, are shown to predict the scalings in the DNS of homogeneous compressible turbulence.

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**2. Preliminaries.** It has been shown, Ristorcelli (1995, 1997), that a low fluctuating Mach number expansion for the compressible Navier Stokes equations produces a diagnostic relationship for the dilatation. The small parameter in these expansions is the square of the fluctuating turbulent Mach number:  $\gamma M_t^2$  and  $M_t = \tilde{u}/c_\infty$ , where  $\tilde{u} = 2k/3 = \langle u_j u_j \rangle / 3$  and  $c_\infty^2 = \gamma P_\infty / \rho_\infty$ . To leading order, the density fluctuations are given, in nondimensional units, by the solenoidal pressure fluctuations,  $\gamma \rho_1 = p_{inc} = p_1$  and the continuity equation becomes a *diagnostic* relation for the fluctuating dilatation,

$$(1) \quad -\gamma d = p_{,t} + v_k p_{,k}.$$

The subscript on  $p_1$  has been dropped. It is seen that one does not need to obtain a solution to the evolution equation for the compressible velocity field,  $w_i$ , in order to obtain its dilatation,  $d = w_{i,i}$ . A very useful result. The dilatation is diagnostically related to the *local* fluctuations of the pressure and velocity; it is the rate of change of the incompressible pressure field  $p_{1,jj} = (v_i v_j)_{,ij}$ , following a fluid particle. Constitutive relations for the pressure-dilatation can be found by taking the appropriate moments of the diagnostic relation for the dilatation to produce,

$$(2) \quad -2\gamma \langle pd \rangle = \frac{D}{Dt} \langle pp \rangle.$$

The near field compressibility effects, as manifested in  $\langle pd \rangle$ , have been directly linked to the solenoidal portions of the velocity field. This fact has been exploited to obtain expressions for the pressure-dilatation covariances, Ristorcelli (1995, 1997):

$$(3) \quad \begin{aligned} \langle pd \rangle = & -\chi_{pd} M_t^2 [P_k - \bar{\rho}\varepsilon + T_k] - \bar{\rho}k M_t^2 \chi_{pd}^r \frac{D}{Dt} \mathcal{T} \\ & + \frac{3}{4} \chi_{pd} M_t^4 \gamma (\gamma - 1) (P_T + \bar{\rho}\varepsilon + T_T) \end{aligned}$$

where

$$(4) \quad \begin{aligned} \chi_{pd} &= \frac{2I_{pd}}{1 + 2I_{pd}M_t^2 + \frac{3}{2}I_{pd}M_t^4\gamma(\gamma-1)} \\ \chi_{pd}^r &= \frac{I_{pd}^r}{1 + 2I_{pd}M_t^2 + \frac{3}{2}I_{pd}M_t^4\gamma(\gamma-1)} \\ I_{pd} &= \frac{2}{3}I_1^s + I_{pd}^r [3\hat{S}^2 + 5\hat{W}^2] \\ I_{pd}^r &= \frac{1}{30} \left(\frac{2}{3}\right)^3 \alpha^2 I_1^r. \end{aligned}$$

Note that  $\mathcal{T} = [3\hat{S}^2 + 5\hat{W}^2]$ . The nondimensional strain and rotation rates are given by:  $\hat{S}^2 = (Sk/\varepsilon_s)^2$ ,  $\hat{W}^2 = (Wk/\varepsilon_s)^2$  where  $S = \sqrt{S_{ij}S_{ij}}$  and  $W = \sqrt{W_{ij}W_{ij}}$ . The strain and rotation tensors are defined in analogy with the incompressible case, *i.e.*, traceless  $S_{ij} = \frac{1}{2}[U_{i,j} + U_{j,i} - \frac{2}{3}D\delta_{ij}]$ ,  $W_{ij} = \frac{1}{2}[U_{i,j} - U_{j,i}]$ . Here  $M_t = \tilde{u}/c_\infty$  where  $\tilde{u} = 2k/3 = \langle u_j u_j \rangle / 3$ ,  $c_\infty^2 = \gamma P_\infty / \rho_\infty$  and  $\alpha$  comes from the Kolmogorov scaling relation  $\ell = \alpha(2k/3)^{3/2}/\varepsilon_s$ . The constants, denoted by the  $I_i$ , in these expressions are given by integrals of the longitudinal correlation function,  $f$ :  $I_1^s = \int_0^\infty \xi f'^2 d\xi$ , and  $I_1^r = 2 \int_0^\infty \xi f d\xi$ . A quick order of magnitude estimate for the integrals can be made using  $f = e^{-\xi^2\pi/4}$ . The following values are found:  $I_1^s = \frac{1}{2}$ ,  $I_1^r = \frac{4}{\pi} = 1.273$ . The values found from high Reynolds number wind tunnel data are similar:  $I_1^s = 0.300$ ,  $I_1^r = 1.392$ , Zhou (1995).

**3. Isotropic decay.** For the isotropic decay the expression for the pressure-dilatation becomes

$$(5) \quad \langle pd \rangle = \chi_{pd} M_t^2 \varepsilon_s .$$

Here  $\bar{\rho}$  has been set to unity. The sign of  $\langle pd \rangle$  is positive indicating a net transfer of energy from potential to kinetic modes. Rearranging, to isolate the scaling, produces

$$(6) \quad \frac{\langle pd \rangle}{M_t^2 \varepsilon_s} = \chi_{pd} = \frac{\frac{4}{3} I_1^s}{1 + \frac{4}{3} I_1^s M_t^2} .$$

Terms of order  $M_t^4$  have been dropped. Earlier estimates given in Ristorcelli (1995), shown above, indicate  $I_1^s = 0.5 - 0.3$ . The theory therefore predicts an asymptote for  $\chi_{pd}$  as the turbulent Mach number vanishes:

$$(7) \quad \chi_{pd} \rightarrow 0.666 - 0.40 \quad \text{as} \quad M_t^2 \rightarrow 0 .$$

The DNS results, shown in Figure 1, were provided by Blaisdell for three different initial turbulent Mach numbers. As a service to the reader, the figure identifies two definitions of the turbulent Mach number: that used by Blaisdell *et al.* (1993) in his simulations,  $M_t^b$ , and that which comes from the perturbation analysis of Ristorcelli (1995, 1997). The present compressible DNS reflect a consistent set of initial conditions as described in Ristorcelli and Blaisdell (1997).

The agreement with the DNS, shown in Figure 1, is very good. The theory has been corroborated without *a posteriori* adjustment of constants. The actual values of the constant could in principle be calculated from the DNS. As they are expected to be weakly dependent on initial conditions, this course is not followed further - what has been presented is sufficient for verification. Moreover, the slow portion of the pressure-dilatation is nominal compared to the rapid portion which is the most important contributor in the shear flows of interest.

**4. Homogeneous shear.** The pressure-dilatation in the homogeneous shear is now investigated. The instantaneous pressure-dilatation is seen in Figure 2. Also shown are its averaged values following the procedure of Sarkar (1992). Here the time integral of the pressure-dilatation has been taken: the vertical axis being  $-\frac{1}{St} \int \frac{\langle pd \rangle}{\varepsilon_s} d(St)$ . The oscillations in the pressure-dilatation associated with the relaxation from initial conditions are not seen. There is, nonetheless, a build up of the oscillations which has been linked to the compressible component of the pressure field, Sarkar (1992), Blaisdell and Sarkar (1993). An explanation consistent with this observed behavior has been advanced in Ristorcelli (1997b).

For homogeneous shear, the expression for the pressure-dilatation can be simplified. For simple shear  $\mathcal{T} = 8\hat{S}^2$  and, neglecting terms of order  $M_t^4$ , one obtains

$$(8) \quad \langle pd \rangle = -\chi_{pd} M_t^2 [P_k - \varepsilon] - k M_t^2 \chi_{pd}^r 8 \frac{D}{Dt} \hat{S}^2 .$$

Here  $\hat{S} = \frac{Sk}{\varepsilon_s}$ . For Blaisdell's homogeneous shear, in which  $S = const.$ , the expression can be rearranged

$$(9) \quad \begin{aligned} \langle pd \rangle = & -\chi_{pd} M_t^2 \varepsilon_s \left[ \frac{P_k}{\varepsilon_s} - 1 \right] \\ & - 16 \chi_{pd}^r M_t^2 \varepsilon_s \left( \frac{Sk}{\varepsilon_s} \right)^2 \frac{D}{D(St)} \frac{Sk}{\varepsilon_s} . \end{aligned}$$

Note that the coefficient of the first term,  $\chi_{pd}$ , ignoring the small slow pressure contribution, scales as  $\chi_{pd} \sim \left(\frac{Sk}{\varepsilon_s}\right)^2$ ; accounting for the definitions of the  $\chi$ 's, the pressure-dilatation can be rewritten as

$$(10) \quad \langle pd \rangle \sim -\left[\alpha^2 \left(\frac{Sk}{\varepsilon_s}\right)^2 M_t^2 \varepsilon_s \left(\frac{P_k}{\varepsilon_s} - 1\right)\right] I_1^r \left[1 + \frac{1}{\frac{P_k}{\varepsilon_s} - 1} \frac{D}{D(St)} \frac{Sk}{\varepsilon_s}\right] .$$

As the flow evolves, it is expected that  $I_1^r \rightarrow const$  and  $[1 + \frac{1}{\frac{P_k}{\varepsilon_s} - 1} \frac{D}{D(St)} \frac{Sk}{\varepsilon_s}] \rightarrow 1$ . The scaling of  $\langle pd \rangle$  with the term in the first set of brackets will be investigated.

**4.1. “Non-equilibrium” aspects of  $\langle pd \rangle$ .** The homogeneous shear DNS is not an equilibrium flow. Using  $\frac{Sk}{\varepsilon_s}$  as an indicator of the non-equilibrium nature of the flow one sees that  $\frac{D}{Dt} \frac{Sk}{\varepsilon_s} \rightarrow 0$  only in the latter stages of the DNS, see Figure 3 bottom plot. Two regions of the flows evolution are accordingly distinguished: a nonequilibrium earlier portion in which  $D/Dt \hat{S} \neq 0$ , and a structural equilibrium portion for which  $D/Dt \hat{S} \rightarrow 0$ . If the flow is in structural equilibrium,  $\frac{Sk}{\varepsilon_s} \rightarrow const$ , one obtains the simplest form of the pressure-dilatation model: denote it by the subscript “*se*”,

$$(11) \quad \langle pd \rangle_{se} = -\chi_{pd} M_t^2 [P_k - \varepsilon].$$

Let this be called the structural equilibrium form of the pressure-dilatation (which does not mean that the flow is in equilibrium as  $P_k \neq \varepsilon$ ).

Figure 3 indicates the relative contributions of the two terms making up the the pressure-dilatation model. It is seen at small times, that both terms make non-negligible contributions to the pressure-dilatation. As the structural equilibrium is approached,  $D/Dt \hat{S} \rightarrow 0$ , and the second term’s importance, as might be expected, becomes negligible. This is manifested in the second graph were the ratio of  $\langle pd \rangle_{se}$  to the total  $\langle pd \rangle$  approaches unity. Also shown is the *percentage* time rate of change of the relative strain  $\frac{\varepsilon_s}{Sk} \frac{D}{Dt} \frac{Sk}{\varepsilon_s}$ . The curves are noisy as they involve differentiation of numerical data; the trends are nonetheless apparent.

**4.2. Pressure-dilatation scalings.** The appropriately scaled integrals of  $\langle pd \rangle$  will now be taken. In this way one can establish whether the scalings predicted by the model are correct. In the latter portions of Blaisdell’s DNS,  $St > 10$ , about three to four eddy turnovers past initialization, a structural equilibrium is approached. In this region the scaling suggested by

$$(12) \quad \langle pd \rangle \sim -[\alpha^2 (\frac{Sk}{\varepsilon_s})^2 M_t^2 \varepsilon_s (\frac{P_k}{\varepsilon_s} - 1)] I_1^r$$

is investigated. After the scaling  $\langle pd \rangle$  the time integral  $\frac{1}{St - St_0} \int ( ) d(St)$  will be taken. Let the symbol  $\int_{ST}$  denote this averaging operation. The integrals

$$\begin{aligned} I_0 &= \int_{ST} \langle pd \rangle \\ I_1 &= \int_{ST} \frac{\langle pd \rangle}{\varepsilon_s} \\ I_2 &= \int_{ST} \frac{\langle pd \rangle}{\varepsilon_s [\frac{P_k}{\varepsilon} - 1]} \\ I_3 &= \int_{ST} \frac{\langle pd \rangle}{M_t^2 \varepsilon_s [\frac{P_k}{\varepsilon} - 1] (\frac{Sk}{\varepsilon})^2} \\ I_4 &= \int_{ST} \frac{\langle pd \rangle}{\alpha^2 M_t^2 \varepsilon_s [\frac{P_k}{\varepsilon} - 1] (\frac{Sk}{\varepsilon})^2} \end{aligned}$$

will be computed from the data. The integrals are shown in Figure 4. The integration starts at  $St = 9$ ; the integrals are normalized by their values at  $St = 10$ . If the scaling suggested by the analysis of Ristorcelli (1995, 1997) is correct the last integral,  $I_4$ , will be approximately constant, reflecting the fact that the time integral  $\frac{1}{St} \int I_1^r d(St) \sim I_1^r$  and  $\frac{D}{Dt} \hat{S} \rightarrow 0$  as the equilibrium portions of the DNS are attained. Inspection of  $I_4$  in Figure 4 shows this to be the case. The period of time  $10 < St < 16$  corresponds to about one eddy turnover time,  $k/\varepsilon$ .

**4.3. The gradient Mach number.** The largest relative collapse of the scaled integrals of  $\langle pd \rangle$  occurs with the quantities  $M_t^2(\frac{Sk}{\varepsilon})^2$  - the collapse from  $I_2$  to  $I_3$ . The quantity  $M_t^2(\frac{Sk}{\varepsilon})^2$  can be interpreted of as the square of the gradient Mach number. The pressure-dilatation is a strong function of the gradient Mach number. Sarkar (1995) has defined a gradient Mach number as  $M_g = S\ell/c$ ; the transverse two-point correlation of the longitudinal velocity is used as the length scale,  $\ell$ . In this article  $\ell$  will be taken as the traditionally defined longitudinal length scale that occurs in the Kolmogorov scaling:  $\ell = \alpha(2k/3)^{3/2}/\varepsilon_s$ . In which case a mean strain Mach number is defined:

$$(13) \quad M_S = \frac{S\ell}{c} = \alpha \frac{2}{3} \frac{Sk}{\varepsilon_s} M_t = \alpha \frac{2}{3} \hat{S} M_t \simeq \frac{2}{3} \hat{S} M_t.$$

In fact, the curve overshoots the optimum collapse (a horizontal line): the gradient Mach number is increasing faster than  $\langle pd \rangle$ .

**4.4. The Kolmogorov scaling coefficient.** The collapse,  $I_4$ , is much better when the Kolmogorov coefficient is included. A new feature associated with compressibility, that of the Kolmogorov scaling coefficient,  $\alpha$ , is thus apparent. The values used for  $\alpha$  come from Blaisdell's DNS. The longitudinal integral length scale,  $\ell$ ,  $k$  and  $\varepsilon_s$  are calculated from the DNS. The Kolmogorov relationship,  $\varepsilon = \alpha(\frac{2}{3}k)^{3/2}/\ell$  is used to find  $\alpha$ .

The pressure-dilatation model is sensitive to variations in the Kolmogorov scaling parameter: in fact it is decreasing as rapidly as  $M_S$  is increasing. The definition for the mean gradient Mach number given above,  $M_S = \frac{2}{3} \frac{Sk}{\varepsilon_s} M_t$ , implicitly assumes  $\alpha \simeq 1$ .  $M_t$  and  $\frac{Sk}{\varepsilon_s}$  are not new quantities for describing turbulence in single point closures;  $\alpha$ , however, is new. This distinction is made in order to recognize  $\alpha$  as a new independent quantity.

The Kolmogorov constant is thought to be a universal constant for high Reynolds number in isotropic turbulence; for non-ideal, finite Reynolds number, anisotropic turbulence  $\alpha$  is a *flow specific quantity*. It is this fact that makes developing a turbulence model from the analysis of Ristorcelli (1995, 1997) difficult: a choice for  $\alpha$  must be made and for any given flow the choice is not, *a priori*, known. The value of  $\alpha$  can be thought of as describing some large scale structural aspects of the flow: it, after all, relates the kinetic energy, its cascade rate and the two-point correlations. These ideas are further developed in Ristorcelli (1997c).

**5. Summary and Conclusions.** The pressure-dilatation is found to be a nonequilibrium phenomena. It scales as  $\alpha^2 M_t^2 (\frac{Sk}{\varepsilon_s})^2 [P_k/\varepsilon_s - 1]$ . For it to be important requires both 1) the square of the gradient Mach number,  $M_S^2$ , to be substantial and 2) for the flow to be out of equilibrium  $P_k \neq \varepsilon$ . The pressure-dilatation has been observed to be either positive or negative. Its dependence on the non-equilibrium nature of the flow,  $P_k \neq \varepsilon$ , indicates that the pressure dilatation can be either a stabilizing or destabilizing. These predictions are consistent with the DNS of Simone *et al.* (1997) who observes such behavior as related to the anisotropy,  $b_{12} = \frac{\langle uv \rangle}{2k}$ .

Except for the well-established Kolmogorov scaling,  $\ell \sim (\frac{2}{3}k)^{3/2}/\varepsilon_s$ , and the quasi-normal assumption, no additional phenomenological assumptions are made. The pressure-dilatation is a function of the Kolmogorov scaling coefficient and this is expected to be an important feature in models for the effects of compressibility on turbulence. The Kolmogorov coefficient is a flow dependent quantity: there is little known about its dependence in non-ideal — anisotropic, strained, inhomogeneous — flow situations. The appearance of the Kolmogorov coefficient, in as much as it links the energy, the spectral flux and a two-point length scale, is an indication of dependence on large scale structure.

In all likelihood, there are many aspects of compressibility that will contribute to the unusual and dramatic stabilizing effects of compressibility. It is unlikely that any one term in a statistical closure will

account fully for the diverse physical mechanisms. This article has focused on one of these effects, the pressure-dilatation. The present analysis treats only the “scalar” effects of compressibility - the reduction  $k$  through the dilatational covariances in the energy budget; it cannot account for the reduction in the shear anisotropy,  $b_{12}$ , or the normal anisotropy,  $b_{22}$ , so important to the production mechanism for the shear stress,  $\langle v_1 v_2 \rangle$ . To account for these more substantial structural effects appears to require a compressible pressure-strain representation accounting for the effects of compressibility. This has been indicated in Blaisdell and Sarkar (1992), Vreman *et al.* (1996) and also Simone *et al.* (1997).

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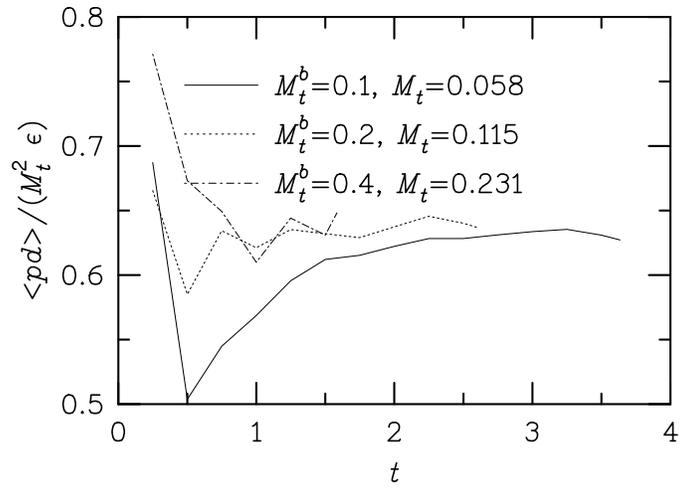


FIG. 1. *The isotropic decay.*

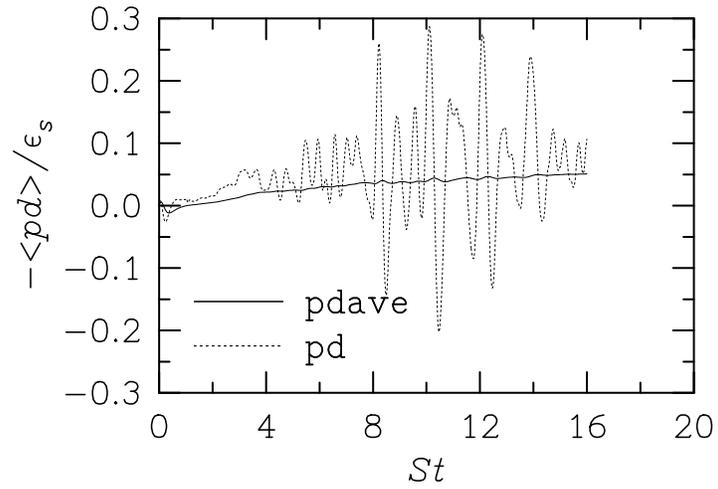


FIG. 2. *The instantaneous and averaged pressure-dilatation - Blaisdell's rjr\_s1 DNS.*

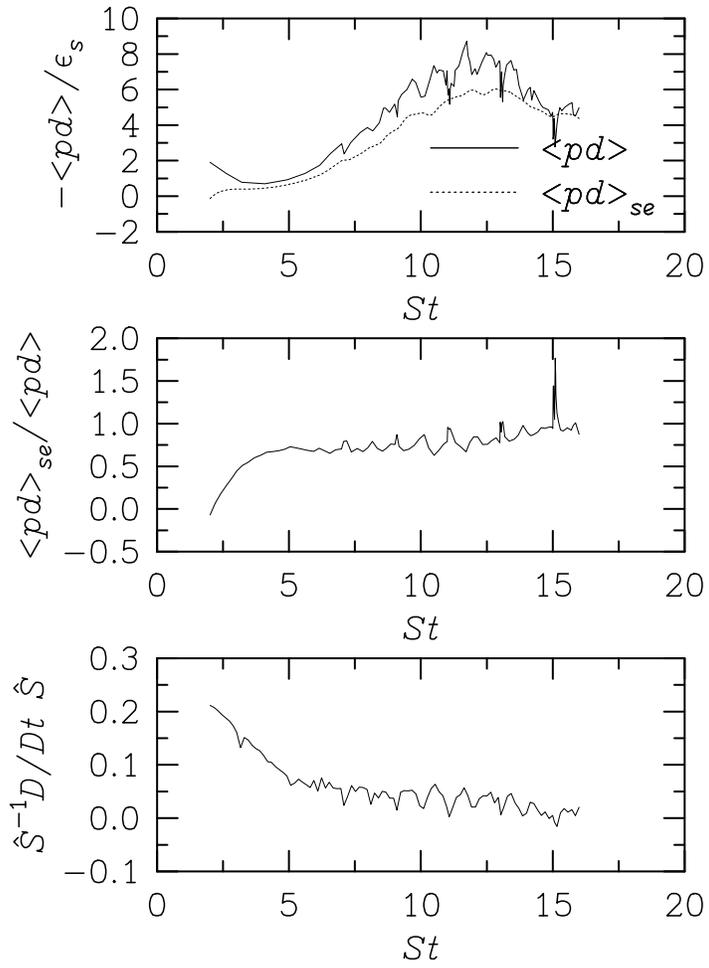


FIG. 3. Contribution of the equilibrium portion of the pressure-dilatation to the pressure-dilatation as predicted by the pseudo-sound analysis.

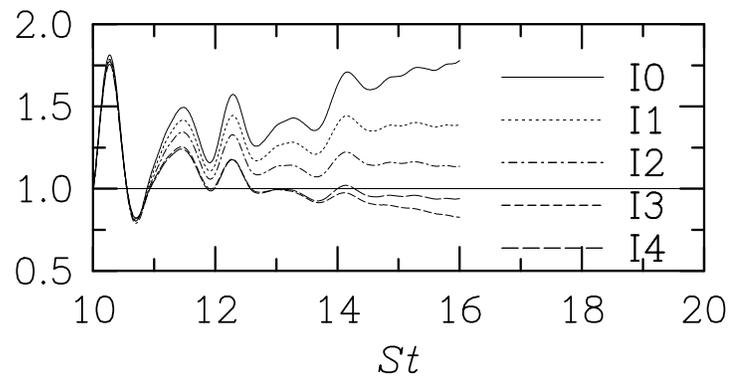
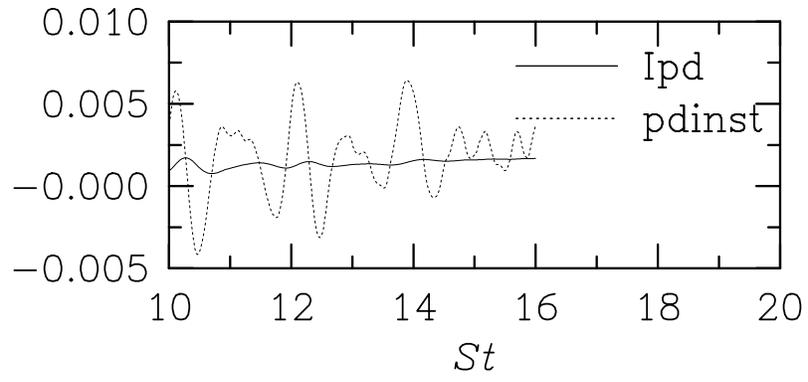


FIG. 4. Integral scalings for the pressure dilatation in homogeneous shear.